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Heavy Quark Colorimetry of QCD Matter

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Abstract

We consider propagation of heavy quarks in QCD matter. Because of large quark mass, the radiative quark energy loss appears to be qualitatively different from that of light quarks at all energies of practical importance. Finite quark mass effects lead to an in-medium enhancement of the heavy–to–light D/π ratio at moderately large (5–10 GeV) transverse momenta. For hot QCD matter a large enhancement is expected, whose magnitude and shape are exponentially sensitive to the density of colour charges in the medium.

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1 Introduction

The energy loss of heavy quarks in QCD matter has excited considerable interest. In particular, collisional energy loss in quark–gluon plasma has been evaluated [1] by the methods of finite–temperature QCD (see [2] for a comprehensive review). However, the energy loss of fast partons in medium is dominated by radiation of gluons [3–5] and, to the best of our knowledge, a consistent treatment of radiative energy loss of heavy quarks in QCD matter is still lacking. Practical importance of this issue has become clear from a number of studies [6] which showed a strong dependence of charmed hadron and lepton spectra on the assumed heavy quark energy loss.

In this letter we follow the recent analysis of the quenching of inclusive particle spectra in QCD medium [7] to compare the yields of hadrons from light and heavy quarks produced in hard interactions in heavy ion collisions. We find that the specific pattern of gluon bremsstrahlung off heavy quarks (the dead cone effect) results in a much smaller heavy quark quenching, leading to a relative enhancement of heavy particle production in heavy ion collisions.

2 Medium induced gluon radiation – a brief summary

We first recall the basic features of gluon radiation caused by propagation of a fast parton (quark) through QCD medium.

As was pointed out in [4], the accompanying radiation is determined by multiple rescattering of the radiated gluon in the medium. The gluon, during its formation time

$$t_{form} \simeq \frac{\omega}{k_\perp^2}\,, \tag{1}$$

accumulates a typical transverse momentum

$$k_\perp^2 \simeq \mu^2 \, \frac{t_{form}}{\lambda}, \tag{2}$$

with λ the mean free path and μ^2 the characteristic momentum transfer squared in a single scattering. This is the random walk pattern with an average number of scatterings given by the ratio t_{form}/λ.

Combining (2) and (1) we obtain

$$N_{coh} = \frac{t_{form}}{\lambda} = \sqrt{\frac{\omega}{\mu^2\,\lambda}} \qquad (3)$$

describing the number of scattering centres which participate, *coherently*, in the emission of the gluon with a given energy ω. For sufficiently large gluon energies, $\omega > \mu^2\lambda$, when the coherent length exceeds the mean free path, $N_{coh} > 1$. In this situation the standard Bethe-Heitler energy spectrum per unit length describing *independent* emission of gluons at each centre gets suppressed:

$$\frac{dW}{d\omega dz} = \left(\frac{dW}{d\omega dz}\right)^{\rm BH} \cdot \frac{1}{N_{coh}} = \frac{\alpha_{\rm s} C_R}{\pi\omega\,\lambda} \cdot \sqrt{\frac{\mu^2\,\lambda}{\omega}} = \frac{\alpha_{\rm s} C_R}{\pi\omega}\sqrt{\frac{\hat{q}}{\omega}}. \qquad (4)$$

Here C_R is the "colour charge" of the parton projectile ($C_R = C_F = \frac{N_c^2-1}{2N_c} = 4/3$ for the quark case we are interested in).

In (4) we have substituted the characteristic ratio μ^2/λ by the so-called gluon *transport coefficient* [8]

$$\hat{q} \equiv \rho \int \frac{d\sigma}{dq^2}\; q^2\; dq^2, \qquad (5)$$

which is proportional to the density ρ of the scattering centres in the medium and describes the typical momentum transfer in the gluon scattering off these centres.

The transport coefficient for cold nuclear matter was expressed in [8] as

$$\hat{q} \simeq \frac{4\pi^2\alpha_s N_c}{N_c^2 - 1}\; \rho\; [xG(x,Q^2)], \qquad (6)$$

with $\rho \simeq 0.16\ {\rm fm}^{-3}$ the average nuclear density and $[xG(x)]$ the gluon density in a nucleon. The latter should be evaluated at the characteristic scale $Q^2 \sim \hat{q}L$, which makes (6), strictly speaking, an equation for $\hat{q}$. Given a slow logarithmic Q^2 dependence of the gluon density, this equation can be solved iteratively. Choosing for the sake of estimate $L \simeq 5$ fm results in $Q \sim 0.5$ GeV. Taking at this scale $\alpha_s \simeq 0.5$ and $[xG(x)] \simeq 1$ (at $x < 0.1$), yields

$$\hat{q}_{\rm cold} \simeq 0.01\ {\rm GeV}^3 \simeq 8\ \rho. \qquad (7)$$

This value is an agreement with the result of the analysis of the gluon $p_\perp$ broadening from the experimental data on J/ψ transverse momentum distributions [9], which in the present notation yielded

$$\hat{q} = (9.4 \pm 0.7)\, \rho\,. \tag{8}$$

An estimate [8] for a hot medium based on perturbative treatment of gluon scattering in quark–gluon plasma with $T \sim 250$ MeV resulted in the value of the gluon transport coefficient of about factor *twenty* larger than (7):

$$\hat{q}_{\rm hot} \simeq 0.2 \ {\rm GeV}^3 \ \simeq \ 20\, \hat{\rm q}_{\rm cold}\,. \tag{9}$$

Multiplying (4) by the length L of the medium traversed,[2] we arrive at the following expression for the inclusive energy distribution of gluons radiated by a quark:

$$\frac{dW}{d\omega} \simeq \frac{\alpha_{\rm s} C_F}{\pi\,\omega}\sqrt{\frac{\omega_1}{\omega}}, \qquad \omega \ < \ \omega_1 \equiv \hat{q}L^2\,. \tag{10}$$

A more accurate treatment shows that the semi-quantitative estimate (4) is valid for $\omega \lesssim \omega_1/2$, while at very large gluon energies $\omega \gg \omega_1$ the gluon yield is very small and decreases as $dI/d\omega \propto \omega^{-3}$ (for details, see [8]).

The fact that the medium induced radiation vanishes for $\omega > \omega_1$ has a simple physical explanation, as according to (3) the formation time of such gluons starts to exceed the length of the medium:

$$t_{form} = \lambda \cdot \sqrt{\frac{\omega}{\mu^2 \lambda}} = \sqrt{\frac{\omega}{\hat{q}}} = L \cdot \sqrt{\frac{\omega}{\omega_1}} \ > \ L\,.$$

Another important feature of medium induced radiation is the relation between the transverse momentum and the energy of the emitted gluon. Indeed, from (1) and (2) (see also (5)) we derive

$$k_\perp^2 \simeq \sqrt{\hat{q}\,\omega}. \tag{11}$$

This means that the angular distribution of gluons with a given energy ω is concentrated at a characteristic energy- (and medium-) dependent emission angle

$$\theta \simeq \frac{k_\perp}{\omega} \sim \left(\frac{\hat{q}}{\omega^3}\right)^{1/4}. \tag{12}$$

[2]For the sake of simplicity we assume here that the medium is static and uniform.

3 Radiation off heavy quarks: dead cone

Gluon bremsstrahlung off a heavy quark differs from the case of a massless parton (produced in a process with the same hardness scale) in one respect: gluon radiation is suppressed at angles smaller than the ratio of the quark mass M to its energy E. Indeed, the distribution of soft gluons radiated by a heavy quark is given by

$$dP = \frac{\alpha_{\mathrm{s}}\, C_F}{\pi}\, \frac{d\omega}{\omega}\, \frac{k_\perp^2\, dk_\perp^2}{(k_\perp^2 + \omega^2\theta_0^2)^2}, \qquad \theta_0 \equiv \frac{M}{E}\,, \tag{13}$$

where the strong coupling constant α_s should be evaluated at the scale determined by the denominator of (13). Equating, in the small-angle approximation, $k_\perp$ with $\omega\theta$ we conclude that the formula (13) differs from the standard bremsstrahlung spectrum

$$dP_0 \;\simeq\; \frac{\alpha_{\mathrm{s}}\, C_F}{\pi}\frac{d\omega}{\omega}\,\frac{dk_\perp^2}{k_\perp^2} \;=\; \frac{\alpha_{\mathrm{s}}\, C_F}{\pi}\frac{d\omega}{\omega}\,\frac{d\theta^2}{\theta^2} \tag{14}$$

by the factor

$$dP_{\mathrm{HQ}} = dP_0 \cdot \left(1 + \frac{\theta_0^2}{\theta^2}\right)^{-2} \tag{15}$$

This effect is known as the "dead cone" phenomenon. Suppression of small-angle radiation has a number of interesting implications, such as perturbative calculability of (and non-perturbative Λ/M corrections to) heavy quark fragmentation functions [10, 11], multiplicity and energy spectra of light particles accompanying hard production of a heavy quark [12].

In the present context we should compare the angular distribution of gluons induced by the quark propagation in the medium with the size of the dead cone. To this end, for the sake of a semi-quantitative estimate, we substitute the characteristic angle (12) into the dead cone suppression factor (15) and combine it with the radiation spectrum (4) to arrive at

$$I(\omega) = \omega\frac{dW}{d\omega} = \frac{\alpha_s\, C_F}{\pi}\sqrt{\frac{\omega_1}{\omega}}\;\frac{1}{(1+(\ell\,\omega)^{3/2})^2}, \tag{16}$$

where

$$\ell \equiv \hat{q}^{-1/3}\,\left(\frac{M}{E}\right)^{4/3}. \tag{17}$$

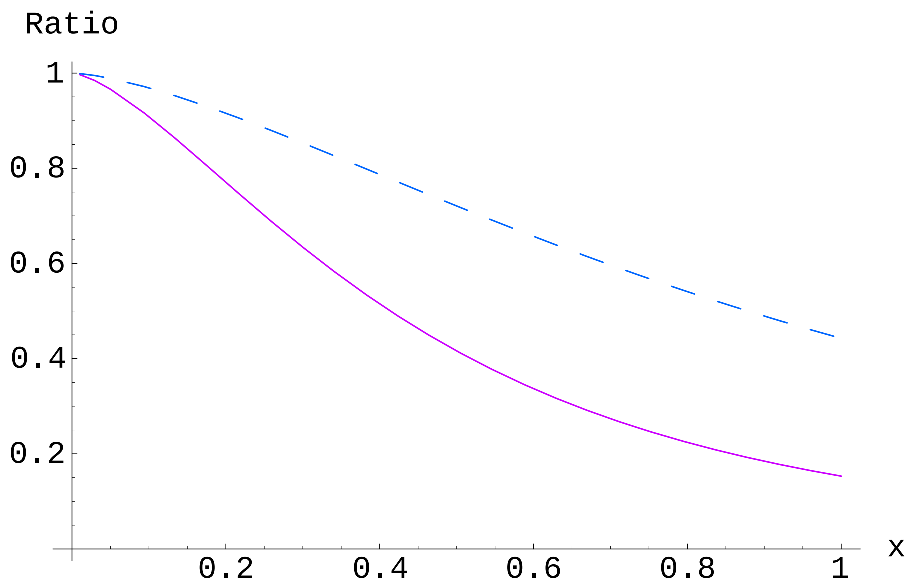

Figure 1: Ratio of gluon emission spectra off charm and light quarks for quark momenta $p_\perp = 10\,\mathrm{GeV}$ (solid line) and $p_\perp = 100\,\mathrm{GeV}$ (dashed); $x = \omega/p_\perp$.

The suppression factor (16) is displayed in Fig. 1 for two quark energies ($x = \omega/p_\perp$).

To see whether the finite quark mass essentially affects the medium induced gluon yield, we need to estimate the product $\ell\omega$ for the maximal gluon energy $\omega \simeq \omega_1$ to which the original distribution (4) extends:

$$\ell\omega_1 = \hat{q}^{-1/3} \left(\frac{M}{E}\right)^{4/3} \cdot \hat{q}L^2 = \left(\frac{E_{\rm HQ}}{E}\right)^{4/3}, \qquad E_{\rm HQ} \equiv M\sqrt{\hat{q}L^3}. \tag{18}$$

This shows that the quark mass becomes irrelevant when the quark energy exceeds the characteristic value $E_{\rm HQ}$ which depends on the size of the medium and on its "scattering power" embodied into the value of the transport coefficient.

Which regime is realized in the experiments on heavy quark production in nuclear collisions? Taking $M = 1.5$ GeV for charm quarks and using the values (7) and (9) we estimate

$$E_{\rm HQ} \;=\; \sqrt{\hat{q}_{\rm cold}}\; L^{3/2}\; M \simeq \; 20\ \mathrm{GeV} \left(\frac{L}{5\,\mathrm{fm}}\right)^{3/2}, \tag{19}$$

$$E_{\mathrm{HQ}} \;=\; \sqrt{\hat{q}_{\mathrm{hot}}}\; L^{3/2}\; M \simeq\; 92\ \mathrm{GeV}\left(\frac{L}{5\,\mathrm{fm}}\right)^{3/2}, \tag{20}$$

for the cold and hot matter, respectively. We observe that for the transverse momentum (energy) distributions of heavy mesons the inequality $E \ll E_{\mathrm{HQ}}$ always holds in practice, especially for the hot medium. We thus conclude that the pattern of medium induced gluon radiation appears to be *qualitatively different for heavy and light quarks* in the kinematical region of practical interest.

4 Quenching

The issue of in-medium quenching of inclusive particle spectra was recently addressed in [7]. The $p_\perp$ spectrum is given by the convolution of the transverse momentum distribution in an elementary hadron–hadron collision, evaluated at a shifted value $p_\perp + \epsilon$, with the distribution $D(\epsilon)$ in the energy ϵ lost by the quark to the medium-induced gluon radiation:

$$\frac{d\sigma^{med}}{dp_\perp^2} = \int d\epsilon\; D(\epsilon)\; \frac{d\sigma^{vac}}{dp_\perp^2}(p_\perp + \epsilon) \equiv \frac{d\sigma^{vac}}{dp_\perp^2}(p_\perp)\cdot Q(p_\perp), \tag{21}$$

with $Q(p_\perp)$ the medium dependent *quenching factor*. The two facts, namely that in the essential region $\epsilon \ll p_\perp$ and that the vacuum cross section is a steeply falling function, allow one to simplify the calculation of the quenching factor Q by adopting the exponential approximation for the ϵ-integral in (21):

$$Q(p_\perp) \simeq \int d\epsilon\; D(\epsilon)\; \exp\left\{\frac{\epsilon}{p_\perp}\cdot\mathcal{L}\right\}, \qquad \mathcal{L} \equiv \frac{d}{d\ln p_\perp}\ln\left[\frac{d\sigma^{vac}}{dp_\perp^2}(p_\perp)\right]. \tag{22}$$

This integral results in the *Mellin moment* of the quark distribution,

$$Q(p_\perp) \;=\; \tilde{D}(\nu) = \exp\left[-\nu\; \int_0^\infty d\omega\; N(\omega)\; e^{-\nu\omega}\right], \quad \nu = \frac{\mathcal{L}}{p_\perp}, \tag{23}$$

where $N(\omega)$ is the *integrated gluon multiplicity* defined according to (see [7] for details)

$$N(\omega) \equiv \int_\omega^\infty d\omega'\, \frac{dW(\omega')}{d\omega'}\,. \tag{24}$$

Here we present an approximate evaluation based on a simplified energy spectrum (16). Integration of (16) leads to the following expression for the gluon multiplicity (24):

$$N(\omega) = \frac{2\alpha_s\, C_F}{\pi}\, \sqrt{\ell\omega_1}\; K(\ell\omega), \tag{25}$$

where

$$\begin{aligned} K(x) \equiv \frac{1}{\sqrt{x}} - \frac{2\pi}{3\sqrt{3}} + \frac{x}{3\,(1+x^{3/2})} + \frac{4}{3\sqrt{3}}\ \arctan\left(\frac{-1+2\sqrt{x}}{\sqrt{3}}\right) \\ + \frac{2}{9}\ \ln\left[\frac{1-\sqrt{x}+x}{(1+\sqrt{x})^2}\right]. \end{aligned} \tag{26}$$

The function $K(x)$ has essential support in the region $\ell\omega = x \ll 1$, that is where the gluon yield is not suppressed by the quark mass (dead cone) effect (cf. (16)). To find the quenching factor we need to evaluate its Mellin transform according to (23), with ν given in (35). The essential energies ω in the multiplicity function in (23) are then restricted to $\omega \lesssim 1/\nu$. It is easy to check that, due to the bias effect (large value of $\mathcal{L} \sim 10$), this restriction is stronger than the quark mass energy cutoff:

$$\frac{\nu}{\ell} = \frac{\mathcal{L}}{p_\perp} \cdot \hat{q}^{1/3} \left(\frac{E}{M}\right)^{4/3} \simeq \mathcal{L} \left(\frac{p_\perp\, \hat{q}}{M^4}\right)^{1/3} \ >\ 2 \tag{27}$$

for $p_\perp > 5$ GeV, even for cold matter. Therefore, in this practically interesting domain, we can approximate

$$N(\omega) \simeq \frac{2\alpha_s\, C_F}{\pi}\ \sqrt{\omega_1}\ \left(\frac{1}{\sqrt{\omega}} - \frac{8\pi\sqrt{\ell}}{9\sqrt{3}}\right) [\, 1 + \mathcal{O}\,(\ell\omega)\,]\,, \tag{28}$$

and get

$$\tilde{D}(\nu) \simeq \exp\left[-\frac{2\alpha_s\, C_F}{\pi}\ \sqrt{\hat{q}}L\ \left(\sqrt{\pi\nu} - \frac{8\pi\sqrt{\ell}}{9\sqrt{3}}\right)\right]. \tag{29}$$

The use of (23) furnishes our final result:

$$Q_H(p_\perp) \simeq \exp\left[-\frac{2\alpha_s C_F}{\sqrt{\pi}}\ L\sqrt{\hat{q}\frac{\mathcal{L}_H}{p_\perp}} + \frac{16\alpha_s C_F}{9\sqrt{3}} L \left(\frac{\hat{q}\ M^2}{M^2+p_\perp^2}\right)^{1/3}\right]. \tag{30}$$

The first term in the exponent in (30) represents the quenching of the transverse momentum spectrum which is universal for the light and heavy quarks,

$$Q_L(p_\perp) \simeq \exp\left[-\frac{2\alpha_s C_F}{\sqrt{\pi}}\ L\sqrt{\hat{q}\frac{\mathcal{L}_L}{p_\perp}}\right]$$

(modulo the difference of the $\mathcal{L}$ parameters determined by the $p_\perp$ distributions in the vacuum). The second term is specific for heavy quarks. It has a positive sign, which means that the suppression of the heavy hadron $p_\perp$ distributions is always smaller than that for the light hadrons. This is a straightforward consequence of the fact that the heavy quark mass suppresses gluon radiation. At very high transverse momenta, both terms vanish – this is in accord with the QCD factorization theorem, stating that the effects of the medium should disappear as $p_\perp \to \infty$. How fast this regime is approached depends, however, on the properties of the medium encoded in the value of the transport coefficient $\hat{q}$ and in the medium size L.

Constructing the ratio of the quenching functions, we estimate the heavy-to-light enhancement factor as

$$\frac{Q_H(p_\perp)}{Q_L(p_\perp)} \simeq \exp\left[\frac{16\alpha_s C_F}{9\sqrt{3}} L\left(\frac{\hat{q}\ M^2}{M^2+p_\perp^2}\right)^{1/3}\right]. \tag{31}$$

This simple expression provides a reasonably good approximation to more accurate quantitative results presented below.

5 Quantitative estimates

To perform a quantitative calculation of the transverse momentum distributions, our considerations have to be combined with a realistic treatment of nuclear geometry, including its centrality dependence, gluon multiple interactions in the initial state (the "Cronin effect"), and a model for the time evolution of QCD matter in the final state. Such an analysis is beyond the scope of the present paper. Nevertheless, we provide some semi-quantitative illustrations of the expected consequences of our results, based on a simplified model of a static uniform medium and a fixed quark path length L.

In what follows we use the full inclusive soft gluon emission spectrum off a massless quark from [8], taking into account the first non-soft correction

$\mathcal{O}(x)$:

$$I(x) \;=\; (1-x)\frac{C_F\alpha_s}{\pi}\ln\left(\cosh^2\left(\sqrt{\frac{\omega_1}{4\omega}}\right)-\sin^2\left(\sqrt{\frac{\omega_1}{4\omega}}\right)\right), \qquad x\equiv\frac{\omega}{E}\,, \quad (32)$$

which, for the heavy quark case, is supplied with the dead-cone suppression factor (15). To illustrate the "dead cone" effect, in Fig. 2 a comparison of the inclusive one-gluon emission spectra off light and charm quarks is shown, for hot medium with $L=5\,\mathrm{fm}$.

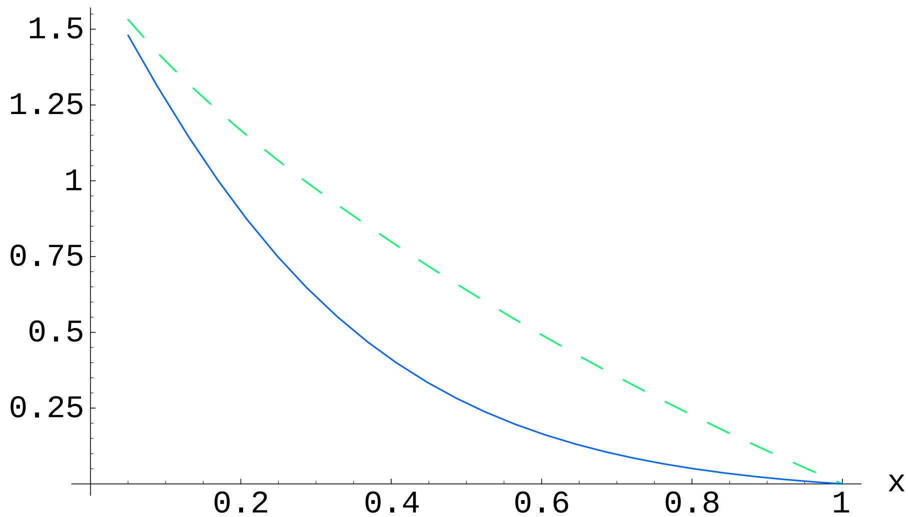

Figure 2: Comparison of energy distributions $\sqrt{x}I(x)$ of gluons radiated off charm (solid line) and light (dashed line) quarks in hot matter with $\hat{q}=0.2\,\mathrm{GeV}^3$ ($p_\perp=10\,\mathrm{GeV}$, $L=5\,\mathrm{fm}$).

As we have discussed above, the quenching of $p_\perp$ distributions of heavy hadrons caused by QCD matter is much weaker than for pions. This is because the gluon cloud around the heavy quark is "truncated" by the large quark mass, in a way which depends on the properties of the medium. This interesting effect can be illustrated by the transverse momentum dependence of the ratio of hadrons originating from the fragmentation of heavy and light quarks (for example, D/π ratio) in heavy ion collisions.

To estimate this ratio we need to know the behaviour of the vacuum spectra. For *light* hadrons (which we assume to originate from light quarks[3])

[3]we remark that the quenching of hadrons produced in the fragmentation of primary

the parameterization of the $p_\perp$ distribution which describes the first RHIC hadroproduction data for $p_\perp$ up to ~ 6 GeV [13]:

$$\frac{d\sigma}{dp_\perp^2} = A \left(\frac{1}{p_0 + p_\perp}\right)^m , \tag{33}$$

with $p_0 = 1.71$ GeV and $m = 12.42$.

The $p_\perp$ distributions of D mesons produced in hadron collisions were experimentally found [14] to be well described by the following simple parameterization:

$$\frac{d\sigma}{dp_\perp^2} = C \left(\frac{1}{bm_c^2 + p_\perp^2}\right)^{\frac{n}{2}} , \tag{34}$$

with $b = 1.4 \pm 0.3$, $n = 10.0 \pm 1.2$, and $m_c = 1.5$ GeV. This parameterization also provides a good fit to the theoretical calculations [15] in perturbative QCD. Using this parameterization in (22) gives

$$\nu \; = \; \frac{n\; p_\perp}{a^2 + p_\perp^2} \, , \qquad a^2 \equiv bm_c^2 \, . \tag{35}$$

for the ν parameter in (23).

Fig. 3 shows the ratio of quenching factors (30) for heavy and light quarks in cold nuclear matter ($L = 5\,\mathrm{fm}$), which is relevant for large-$p_\perp$ particle production in pA collisions. A small value of the transport coefficient, $\hat{q}_{\rm cold} \simeq 0.01\,\mathrm{GeV}^3$, translates into a $\sim 15\%$ enhancement of the heavy-to-light ratio (about 1% for $L = 2\,\mathrm{fm}$).

According to (31), we expect a larger quenching ratio for a hot medium. Indeed, as Fig. 4 demonstrates, the D/π ratio should become significantly enhanced as compared to pp collisions: assuming a fixed length $L = 5\,\mathrm{fm}$ of the hot medium traversed by the quarks, we find a factor of ~ 2 enhancement at $p_\perp \sim 5 \div 10$ GeV.

We only present the *ratio* of quenching factors for the following reason. As noticed in [7], in the region of transverse momenta we are considering, the *absolute* magnitude of quenching turns out to be extremely sensitive to gluon radiation in the few-hundred-MeV energy range and therefore cannot be quantitatively predicted without detailed understanding of the spectral properties of the medium.

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gluons contains an additional factor $9/4 \simeq 2$ in the exponent

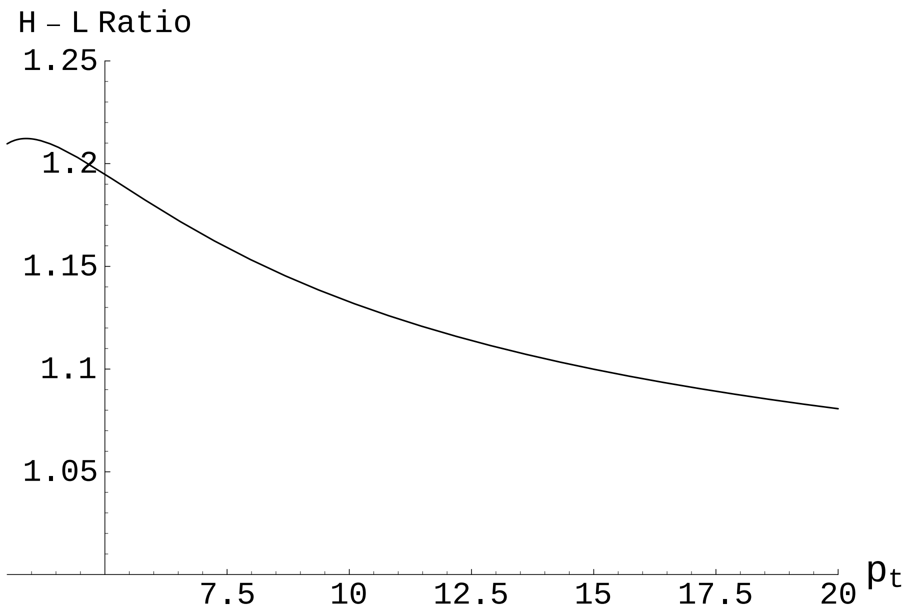

Figure 3: The ratio of quenching factors $Q_H(p_\perp)/Q_L(p_\perp)$ for charm and light quarks in cold nuclear matter ($\hat{q} = 0.01\,\mathrm{GeV}^3$, $L = 5\,\mathrm{fm}$).

The heavy-to-light ratio, however, proves to be much less sensitive to the infrared region, since gluon radiation off heavy and light quarks is universal in the $x \to 0$ limit, see (16). To illustrate this point, in Fig. 4 we show the ratio of quenching factors calculated with the energy restriction upon gluon energies, $\omega > 500\,\mathrm{MeV}$. We see that in the $5 - 10\,\mathrm{GeV}$ range of $p_\perp$, this modifies the ratio by $20 - 30$ %.[4]

The predicted ratio should not be taken at its face value: the enhancement factor originates mainly from a very strong quenching of light quark jets, and in reality such a dramatic suppression will be washed out by light-particle production at the *periphery* of the collision zone. In a central AA collision, the number of collisions in this peripheral zone, depending on the strength of absorption, scales as $\sim A^{1/3} \div A^{2/3}$, whereas the total number of collisions

[4]The quenching factors themselves change (increase) by an order of magnitude when the radiation of gluons with energies smaller than $500\,\mathrm{MeV}$ is vetoed.

which contribute to the production of D mesons scales as $\sim A^{4/3}$. Thus, even with account of the collision geometry, we may still expect an order of $\sim A^{2/3} \div A$ enhancement of the D/π ratio, which for a $Au - Au$ collision may translate in a strong effect.

Moreover, given a large nucleus and moderately large $p_\perp$, both light and (especially) heavy quarks will be tempted to turn into hadrons *prior* to leaving the interaction arena. In the pion case, this is likely to add to quenching via pion absorption in the medium [16], thus enhancing further the D/π ratio.

Clearly, detailed calculations have to be performed before a defendable number for the magnitude of the D/π enhancement can be presented. Nevertheless, the D/π ratio appears to be extremely sensitive to the density of colour charges in QCD matter, and we eagerly await the results of experimental studies of this quantity in relativistic heavy ion collisions.

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References

[1] E. Braaten and M. Thoma, Phys. Rev. D44 (1991) R2625;
M. Thoma and M. Gyulassy, Nucl. Phys. B351 (1991) 491;
R. Baier, M. Dirks and K. Redlich, hep-ph/9910353.

[2] J.-P. Blaizot and E. Iancu, hep-ph/0101103.

[3] M. Gyulassy and X.-N. Wang, Nucl. Phys. B420 (1994) 583.

[4] R. Baier, Yu.L. Dokshitzer, S. Peigne and D. Schiff, Phys.Lett. B345 (1995) 277.

[5] B.G. Zakharov, JETP Lett. 63 (1996) 952.

[6] E. Shuryak, Phys. Rev. C55 (1997) 961;
Z. Lin, R. Vogt and X.-N. Wang, Phys. Rev. C57 (1998) 899;
M.G. Mustafa, D. Pal, D.K. Srivastava and M. Thoma, Phys. Lett. B428 (1998) 234.

[7] R. Baier, Yu.L. Dokshitzer, A.H. Mueller, and D. Schiff, preprint BI-TP 2001/02, LPT-Orsay-01/57, under preparation.

[8] R. Baier, Yu.L. Dokshitzer, A.H. Mueller, S. Peigne and D. Schiff, Nucl.Phys. B484 (1997) 265.

[9] D. Kharzeev, M. Nardi and H. Satz, Phys. Lett. B405 (1997) 14.

[10] Yu.L. Dokshitzer, V.A. Khoze and S.I. Troyan, Phys.Rev.D53 (1996) 89-119 [hep-ph/9506425]

[11] P. Nason and B.R. Webber, Phys.Lett. B395 (1997) 355 [hep-ph/9612353]

[12] B.A. Schumm, Yu.L. Dokshitzer, V.A. Khoze and D.S. Koetke, Phys.Rev.Lett. 69 (1992)3025;
Yu.L. Dokshitzer, V.A. Khoze and S.I. Troyan, J.Phys.G17 (1991) 1481.

[13] G. David, for the PHENIX Coll., *contribution to the Quark Matter 2001 Conf.* [nucl-ex/0105014]
J. Velkovska, for the PHENIX Coll., *contribution to the Quark Matter 2001 Conf.* [nucl-ex/0105012]
C. Adler et al., the STAR Coll., [nucl-ex/0106004]
J.C. Dunlop, for the STAR Coll., *contribution to the Quark Matter 2001 Conf.* [http://www.rhic.bnl.gov/qm2001].

[14] C.A. Alves et al, The E769 Coll., Phys. Rev. Lett. 77 (1996) 2392.

[15] S. Frixione, M.L. Mangano, P. Nason and G. Ridolfi, hep-ph/9702287, and references therein.

[16] B.Z. Kopeliovich, J. Nemchik and E. Predazzi, in: Proceedings of Workshop on Future Physics at HERA, Hamburg, Germany, 30-31 May 1996 [nucl-th/9607036].

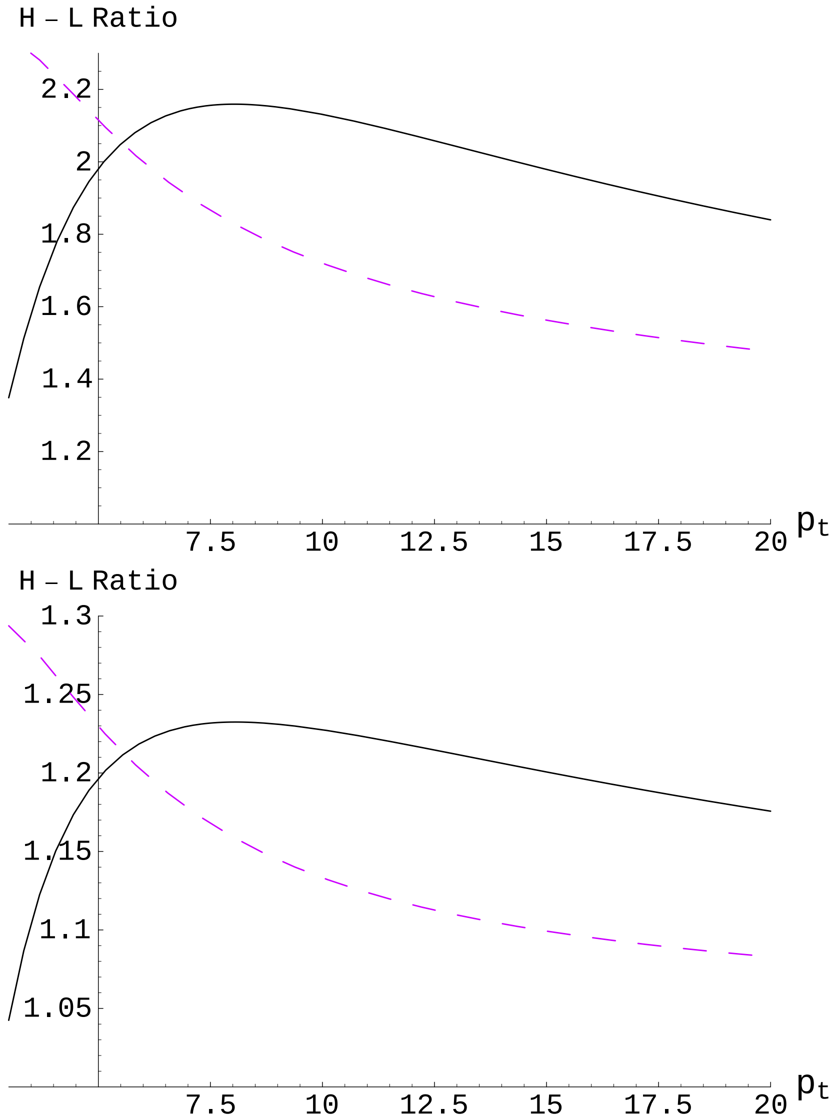

Figure 4: The ratio of quenching factors $Q_H(p_\perp)/Q_L(p_\perp)$ for charm and light quarks in hot matter with $\hat{q} = 0.2\,\mathrm{GeV}^3$ ($L = 5\,\mathrm{fm}$ upper panel; $L = 2\,\mathrm{fm}$ lower panel). Solid lines correspond to unrestricted gluon radiation, while the dashed lines are based on the calculation with the cut on gluon energies $\omega > 0.5\,\mathrm{GeV}$.